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Two Higgs Bi-doublet Left-Right Model With Spontaneous P and CP Violation

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Abstract

A left-right symmetric model with two Higgs bi-doublet is shown to be a consistent model for both spontaneous P and CP violation. The flavor changing neutral currents can be suppressed by the mechanism of approximate global $U(1)$ family symmetry. We calculate the constraints from neural K meson mass difference Δm_K and demonstrate that a right-handed gauge boson W_2 contribution in box-diagrams with mass well below 1 TeV is allowed due to a cancellation caused by a light charged Higgs boson with a mass range $150 \sim 300$ GeV. The W_2 contribution to ϵ_K can be suppressed from appropriate choice of additional CP phases appearing in the right-handed Cabbibo-Kobayashi-Maskawa matrix. The model is also found to be fully consistent with B^0 mass difference Δm_B, and the mixing-induced CP violation quantity $\sin 2\beta_{J/\psi}$, which is usually difficult for the model with only one Higgs bi-doublet. The new physics beyond the standard model can be directly searched at the colliders LHC and ILC.

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I. INTRODUCTION

Since the discovery of parity (P) violation fifty years ago[1, 2], it has been realized that both symmetry and asymmetry can play important roles in particle physics. Later on, charge-conjugation-parity (CP) violation was also discovered in kaon decays[3]. The electroweak standard model was established based on the left-handed symmetry $SU(2)_L$ and has well been described by the gauge symmetry $SU(2)_L \times U(1)_Y$[4, 5, 6]. Since then, one of the important issues in particle physics concerns origin of P and CP violations as well as the smallness of flavor changing neutral currents(FCNC). Its solution requires physics beyond the standard model.

The investigation of explicit CP violation in the standard model led to the prediction for the existence of three generation quarks, so that a single Kobayashi-Maskawa CP phase[7] can be introduced to characterize the CP-violating mechanism in the standard model. Such a simple CP-violating mechanism has been found to be remarkable for explaining not only the indirect CP violation in kaon decays, but also the direct CP violation in kaon decays[8, 9] observed by two experimental groups at CERN[10] and Fermilab[11], and the direct CP violations in B meson decays[12, 13, 14] reported by two B-factories[15, 16]. Nevertheless, the CP violation in the standard model is assumed to be caused from the explicit complex Yukawa couplings put in by hand, thus its origin remains unknown. To understand the origin of CP violation, a spontaneous CP violation mechanism was suggested by Lee in 1973[17, 18] in which scalar fields are responsible to CP violation. Soon after, an interesting spontaneous CP-violating three Higgs doublet model was proposed by imposing discrete symmetries[19] in order to avoid the FCNC, while such a model has been strongly constrained from the low energy phenomena of K and B systems. By abandoning the natural flavor conservation hypothesis[20, 21, 22], a general two Higgs doublet model (2HDM) motivated from spontaneous CP violation has been investigated in detail[21, 22], where the FCNC is assumed to be naturally suppressed by the mechanism of approximate global U(1) family symmetry[21, 22]. Of particular, it has been shown in refs.[21, 22] that after spontaneous symmetry breaking the single relative CP phase of two vacuum expectation values can induce rich CP-violating sources, which not only explain the KM CP-violating mechanism in the standard model, but also lead to new type of CP-violating sources in the charged Higgs interactions. Such a model can result in new physics phenomena[23, 24, 25, 26, 27] and

remains consistent with the current experiments.

With the hypothesis that parity is a good symmetry at high energy, a left-right symmetric model was proposed based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$[28, 29, 30]. In such a model, parity violation can naturally be understood via spontaneous symmetry breaking. Also CP asymmetry can be realized as a consequence of spontaneous symmetry breaking[31, 32, 33, 34]. Nevertheless, the spontaneous P and CP-violating left-right model with only one Higgs bi-doublet is strongly constrained[32, 33, 34] from low energy phenomenology:

(i) The neutral kaon mass difference Δm_K requires that the right-handed gauge bosons must be very heavy above 2 TeV to suppress the extra box-diagram as the gauge coupling is left-right symmetric. Since the Yukawa couplings for neutral and charged Higgs bosons are fixed to quarks masses and Cabbibo-Kobayashi-Maskawa (CKM) matrices. There is no cancellation occurring among different contributions. For the same reason, the lightest neutral Higgs boson must be above 10 TeV[34, 35, 36] to suppress FCNC. Such a neural Higgs mass is too heavy in the Higgs bi-doublet sector to make the model natural as the bi-doublet Higgs bosons are expected to be at the electroweak scale which is much lower than the right-handed gauge boson mass; (ii) In the one Higgs bi-doublet model, all the CP violating phases are calculable quantities in terms of quarks masses and ratios of VEVs, which can be directly tested by the experimental data on CP violating observables. It has been shown [34] that the combing constraints from K system and B system actually excluded the so-called minimal one Higgs bi-doublet left-right model with spontaneous CP violation in the decoupling limit, as the model fails to reproduce the precisely measured weak phase angle $\sin 2\beta$ from B factories; (iii) Furthermore, the condition for the spontaneous CP violation requires an unnatural fine tuning of the Higgs potential in the one Higgs bi-doublet LR model[36, 37]. For those reasons, it was motivated to consider the one Higgs bi-doublet LR model with general CP violation[38, 39, 40] instead of spontaneous CP violation. An alternative consideration for spontaneous P and CP violation is to introduce the concept of mirror particles[41, 42], recently, a maximally symmetric model[43, 44] was constructed along this line by considering mirror quarks and leptons.

In this note, motivated by the general 2HDM as a model for spontaneous CP violation, we shall simply extend the one Higgs bi-doublet left-right model to a two Higgs bi-doublet left-right model with spontaneous P and CP violation, and demonstrate that the above

mentioned stringent phenomenological constraints from neutral meson mixings can be significantly relaxed. It will be shown that the right-handed gauge boson mass can be as low as 600 GeV with the charged Higgs mass around 200 GeV. The FCNC will not impose severe constraints on the neural Higgs mass, provided small off-diagonal Yukawa couplings via the mechanism of approximate global $U(1)$ family symmetry[20, 21, 22].

The paper is organized as follows: in section II, we present a general description for two Higgs bi-doublet left-right model. In section III, we analyze the neutral K system, which includes the mass difference Δm_K and indirect CP violation ϵ_K. We observe that the right-handed gauge boson contributions to the mass difference Δm_K can be opposite to that from the charged Higgs boson in this extended model and a cancellation between the two contributions is possible in a large parameter space. The suppression of right-handed gauge boson contributions to the indirect CP violation ϵ_K is found to occur naturally. As a consequence, a light right-handed gauge boson around the current experimental low bound is allowed. In section IV, we discuss in detail the neutral B meson system, the mass difference Δm_B and the time dependent CP asymmetry in $B^0 \to J/\Psi K_S$ decay are found to be consistently characterized in the two Higgs bi-doublet model with spontaneous P and CP violation, which is unlike the one Higgs bi-doublet model. Conclusions and remarks are presented in the last section.

II. PROPERTIES OF A TWO HIGGS BI-DOUBLET LEFT-RIGHT MODEL

The left-handed and right-handed quarks and leptons in the $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ model are all given by the doublets

$$\begin{aligned} Q_{iL} &= \begin{pmatrix} u_i \\ d_i \end{pmatrix}_L : (2,1,1/3), \quad Q_{iR} = \begin{pmatrix} u_i \\ d_i \end{pmatrix}_R : \quad (1,2,1/3), \\ L_{iL} &= \begin{pmatrix} \nu_i \\ l_i \end{pmatrix}_L : \quad (2,1,-1), \quad L_{iR} = \begin{pmatrix} \nu_i \\ l_i \end{pmatrix}_R : \quad (1,2,-1), \end{aligned} \tag{1}$$

where $i = 1,2,3$ runs over number of generations. The quantum numbers (X_L, X_R, Y) in parenthesis denote the $SU(2)_L$, $SU(2)_R$ and $U(1)_{B-L}$ representation. $X_{L,R}$ represent dimensions of the $SU(2)_L$ and $SU(2)_R$ representations, and Y is the hypercharge $Y = B-L$.

As a gauge invariant model, three gauge fields for the symmetry group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ are introduced as W_L^μ, W_R^μ and B^μ respectively. The gauge invariant fermion-gauge interactions are constructed as follows

$$\mathcal{L}_f = \sum_{\Psi=(Q),(L)} \bar{\Psi}_L \gamma^\mu \left(i\partial_\mu + g_L \frac{\tau^i}{2} W_{L\mu}^i + g' \frac{Y}{2} B_\mu \right) \Psi_L + (L \to R) \,. \tag{2}$$

To generate masses of fermions and gauge bosons, we shall introduce scalar fields and apply the Higgs mechanism to break symmetry spontaneously. In order to generate fermion mass matrices, one only needs to introduce one Higgs bi-doublet [29, 30]. However, in view of the above mentioned phenomenological difficulties, here we shall consider a left-right symmetric model with two Higgs bi-doublets

$$\phi = \begin{pmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{pmatrix}, \quad \chi = \begin{pmatrix} \chi_1^0 & \chi_1^+ \\ \chi_2^- & \chi_2^0 \end{pmatrix} \quad : (2,2,0) \,. \tag{3}$$

The most general Yukawa interaction for quarks is given by

$$\mathcal{L}_Y \;=\; -\sum_{i,j} \bar{Q}_{iL} \left((y_q)_{ij}\phi + (\tilde{y}_q)_{ij}\tilde{\phi} + (h_q)_{ij}\chi + (\tilde{h}_q)_{ij}\tilde{\chi} \right) Q_{jR}, \tag{4}$$

where $\tilde{\phi}(\tilde{\chi}) = \tau_2 \phi^*(\chi^*)\tau_2$ also belong to the representation $(2,2,0)$. Parity P symmetry requires $g_L = g_R \equiv g$ and

$$y_q \;=\; y_q^\dagger, \quad \tilde{y}_q = \tilde{y}_q^\dagger, \quad h_q = h_q^\dagger, \quad \tilde{h}_q = \tilde{h}_q^\dagger. \tag{5}$$

When both P and CP are required to be broken down spontaneously, all the Yukawa coupling matrices are real symmetric.

Note that allowing the two Higgs bi-doublet coupling to the same quark field may generate large FCNC at tree level. To suppress FCNC, we shall follow the similar treatment in the general two-Higgs-doublet model[21, 22] by considering the mechanism of approximate global $U(1)$ family symmetry[20, 21, 22]

$$(u_i, d_i) \to e^{-i\theta_i}(u_i, d_i), \tag{6}$$

which is motivated by the approximate unity of the CKM matrix. As an consequence, y, $\tilde{y}$, h and $\tilde{h}$ are nearly diagonal matrices.

To break $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ to the $U(1)_{em}$, it requires to introduce other Higgs multiplets in addition to ϕ and χ. The most popular choice for generating small neutrino masses is to introduce Higgs triplets ($\Delta_L \sim (3,1,2)$, $\Delta_R \sim (1,3,2)$) [29, 30]

$$\Delta_{L,R} = \begin{pmatrix} \delta^+_{L,R}/\sqrt{2} & \delta^{++}_{L,R} \\ \delta^0_{L,R} & -\delta^+_{L,R}/\sqrt{2} \end{pmatrix} e . \tag{7}$$

Introducing the SU(2) triplets breaks the custodial symmetry and leads to corrections to the parameter $\rho = m_W^2/(m_Z^2 \cos^2\theta_W)$ from unity at tree level, which may subject to strong constraints from the LEP data. However, the corrections from the model involve a number of free parameters such as the ratios among the VEVs of Higgs bidoublets and triplets. The constraint from a single ρ parameter is not severe. Furthermore, the low energy process such as μ decays and neutral current interactions νN , νe and eN are all affected by the model parameters, which modifies the SM relations among the electroweak precision observables. Useful constraints can be obtained from a global fit to all the relevant data. It has been shown that the combined analysis on both the high and low energy electroweak data within the minimal left-right model only leads to a mild lower bound of $M_2 > 700 \sim 800$ GeV for right-hande gauge boson W_2 when the correction to ρ parameter is below 1%. Since the two Higgs bidoublet model considered here contains one more bidoublet which does not violate the custodia symmetry and contain more parameters, the constraints should be even weaker.

The most general form of Higgs potential in this model is rather complicated, which involves the quadratic and quartic terms for the extra bi-doublet field χ and its mixing with ϕ and the two triplets $\Delta_{L,R}$. It has been shown that by simply adding a singlet Higgs field to the one Higgs bi-doublet LR model, the spontaneous CP violation can occur naturally[45]. The two Higgs bi-doublet model has definitely more flexibility in Higgs potential. It is expected that in the most general case the spontaneous CP violation is allowed, which will not be discussed in detail in the present note.

After the spontaneous symmetry breaking, the two Higgs bi-doublet fields can have the following vacuum expectation values(VEVs)

$$\langle\phi\rangle = \begin{pmatrix} v_1 e^{i\delta_1} & 0 \\ 0 & v_2 e^{i\delta_2} \end{pmatrix} \quad \text{and} \quad \langle\chi\rangle = \begin{pmatrix} w_1 e^{i\varphi_1} & 0 \\ 0 & w_2 e^{i\varphi_2} \end{pmatrix} , \tag{8}$$

which leads to the following mass matrix for quarks

$$
\begin{aligned}
M_u &= y_q v_1 e^{i\delta_1} + \tilde{y}_q v_2 e^{-i\delta_2} + h_q w_1 e^{i\varphi_1} + \tilde{h}_q w_2 e^{-i\varphi_2}, \\
M_d &= y_q v_2 e^{i\delta_2} + \tilde{y}_q v_1 e^{-i\delta_1} + h_q w_2 e^{i\varphi_2} + \tilde{h}_q w_1 e^{-i\varphi_1}.
\end{aligned} \tag{9}
$$

As the mass matrices are symmetric, they can be diagonalized by

$$
M_u = U M_u^D U^T \quad \text{and} \quad M_d = V M_d^D V^T, \tag{10}
$$

with $M^D_{u(d)}$ being diagonal mass matrices for up(down)-type quarks. It follows that the resultant quark mixing matrices for left-handed and right-handed quarks are complex conjugate to each other

$$
K^L = U^\dagger V \quad \text{and} \quad K^R = U^T V^* = K_L^*. \tag{11}
$$

Note that rotating the left-handed quark mixing matrix to the standard CKM form V^L with a single CP phase is non-trivial due to the existence of right-handed quark-gauge interactions and charged Higgs Yukawa interactions. In this case, there are in general five additional CP phases $\alpha_i, (i = 1, 2, 3)$ and $\beta_i, (i = 1, 2)$ in the right-handed quark mixing matrix which is parametrized as follows

$$
V^R = \eta^u \begin{pmatrix}
(V^L_{ud})^* e^{2i\alpha_1} & (V^L_{us})^* e^{i(\alpha_1+\alpha_2+\beta_1)} & (V^L_{ub})^* e^{i(\alpha_1+\alpha_3+\beta_1+\beta_2)} \\
(V^L_{cd})^* e^{i(\alpha_1+\alpha_2-\beta_1)} & (V^L_{cs})^* e^{2i\alpha_2} & (V^L_{cb})^* e^{i(\alpha_2+\alpha_3+\beta_2)} \\
(V^L_{td})^* e^{i(\alpha_1+\alpha_3-\beta_1-\beta_2)} & (V^L_{ts})^* e^{i(\alpha_2+\alpha_3-\beta_2)} & (V^L_{tb})^* e^{2i\alpha_3}
\end{pmatrix} \eta^d. \tag{12}
$$

The sign matrices $\eta^{u,d}$ corresponds to the 32 different sign arrangements of quark masses[34]. In the following considerations, we should focus only on the case with positive quark masses, i.e. $\eta^u = \eta^b = \mathbb{1}$.

Within the Wolfenstein parametrization, we define $\beta_L \equiv \arg(V^{L*}_{td} V^L_{tb})$ which is, to a high precision, one of the angles of the unitarity triangle. As V^L_{tb} is real by convention, one has $\beta_L = \arg(V^{L*}_{td})$. Under this parametrization, one can define a similar quantity $\beta_R \equiv \arg(V^{R*}_{td} V^R_{tb})$. They satisfy the following relation

$$
\beta_L + \beta_R = -(\alpha_1 - \alpha_3 - \beta_1 - \beta_2), \tag{13}
$$

which is useful for discussing B meson system. Similarly, one can define phase parameters relevant to the K meson system, i.e., $\beta'_L \equiv \arg(V^{L*}_{td} V^L_{ts})$ and $\beta'_R \equiv \arg(V^{R*}_{td} V^R_{ts})$. They are

related by

$$\beta_L' + \beta_R' = -(\alpha_1 - \alpha_2 - \beta_1). \tag{14}$$

In the Wolfenstein parametrization, it is known that $|\arg(V_{ts}^L)/\arg(V_{td}^L)| \ll 1$, we have in a good approximation

$$\beta_L' \simeq \beta_L \quad \text{and} \quad \beta_R' = \beta_R + \alpha_2 - \alpha_3 - \beta_2. \tag{15}$$

With enlarged Higgs sector, in this model there are two doubly charged Higgs $H_i^{++}, (i = 1, 2)$, four singly charged Higgs particles $H_i^+, (i = 1, \dots 4)$, six neutral scalars $h_i^0, (i = 1, \dots 6)$ and four neural pseudo-scalars $A_i^0, (i = 1, \dots 4)$. The doubly charged H_i^{++} contribute only to the leptonic sector such as lepton flavor violation processes[46], whereas the singly charged and neutral scalars may have significant effects on mixings and CP violation in quark sector. For the sake of simplicity, we should work in a simplest scenario that only one charged Higgs (labeled as H^+) is light enough to actively contribute to the box-diagrams. Of particular, when the VEVs satisfy the conditions $v_2 \ll v_1$ and $w_2 \ll w_1$ or more precisely $v_2/v_1, w_2/w_1 < m_b/m_t$, which are also needed for obtaining the experimentally allowed small mixing between left-handed and right-handed gauge bosons, then many features of this model are similar to the general 2HDM with spontaneous CP violation[21, 22]. Therefore, we consider here the 2HDM-like charged Higgs to be the lightest one, the corresponding Yukawa interaction is parametrized as follows

$$\mathcal{L}_C \; = \; -(2\sqrt{2}G_F)^{1/2}\bar{u}^i \left(\sqrt{m_i^u m_k^u} \; \xi_{ik}^u V_{kj}^L P_L - V_{ik}^L \sqrt{m_k^d m_j^d} \; \xi_{kj}^d P_R \right) d^j \; H^+ + \text{H.c.} \tag{16}$$

Here we have used the Cheng-Sher parametrization[47] in the general 2HDM with $\xi_{ij}^{u(d)}$ the effective Yukawa coupling matrices in the physics basis after spontaneous symmetry breaking. The small off-diagonal terms characterized in $\sqrt{m_i^q m_j^q} \; \xi_{ij}^q$ describe the breaking of the global $U(1)$ family symmetry. For future convenience, we denote the diagonal elements as $\xi_c \equiv \xi_{22}^u$ and $\xi_t \equiv \xi_{33}^u$ etc.

For the flavor changing neutral Higgs boson interactions, in the same case that $v_2/v_1, w_2/w_1 < m_b/m_t$, it becomes similar to the general 2HDM with spontaneous CP violation. When considering the 2HDM-like neutral Higgs boson h^0 to be the lightest one, the dominant interactions can approximately be expressed in the following form

$$\mathcal{L}_N = -(\sqrt{2}G_F)^{1/2}\bar{q}_L^i \sqrt{m_i^q m_j^q} \; \eta_{ij}^q \; q_R^j \; h^0 + \text{H.c} \tag{17}$$

where η^q_{ij} are given by ξ^q_{ij} up to the factors caused by the mixing matrix elements O_{ij} among the neutral Higgs bosons, i.e., $\eta^q_{ij} \sim (O_{1k} \pm iO_{1l})\xi^q_{ij}$. Note that a remarkable difference from the one Higgs bi-doublet LR model with spontaneous CP violation is that the effective Yukawa couplings ξ^q_{ij} or η^q_{ij} in the physics basis are in general all complex and no longer symmetric due to spontaneous P and CP violation, they contain more free parameters due to the extra source of CP violation in the VEVs and more Yukawa couplings associated with the extra Higgs bi-doublet. As a consequence, the effective Yukawa couplings η^q_{ij} or ξ^q_{ij} and the CKM matrices V^L and V^R are no longer directly linked to the quark masses and the ratios of VEVs. Similar to the two Higgs doublet model with spontaneous CP violation[21, 22], the effective Yukawa couplings η^q_{ij} or ξ^q_{ij} in the two Higgs bi-doublet model are also free parameters which can cause significantly different effects in low energy phenomenology.

The neutral meson mixing can arise from the neutral scalar exchange at tree level, which could be significant. The contributions to the mixing matrix between the neutral meson P^0 and $\bar{P}^0$ can easily be obtained. Denoting P^0 the bound state of two quarks with quantum number $P^0 \equiv (\bar{q}_i\gamma_5 q_j)$, we have, in the factorization approximation, the following general form

$$M_{12} = \frac{1}{2m_P}\langle P^0|H_{eff}|\bar{P}^0\rangle \simeq G_F \frac{\sqrt{2} f_P^2 m_P B_P^S}{4m^2_{h^0}} \left[\frac{1}{6} + \frac{m_i^q m_j^q}{\left(m_i^q + m_j^q\right)^2}\right] (\eta^q_{ij} - \eta^{q*}_{ji})^2. \tag{18}$$

III. NEUTRAL K MESON MIXING

We proceed to discuss the low energy phenomenological constraints of this model. Since $W_L W_R$ mixing angle is very small from μ decays, the mass eigenstate $W_{1(2)}$ is almost left(right)-handed. In the left-right model, the strongest constraint comes from K meson system. The K^0 meson receives additional contributions from both $W_1 W_2$ loop and charged Higgs loop in box-diagrams. As the internal (c, c) quark loop dominates the whole contribution, any CP violating phases associated with it have to be strong suppressed in order to accommodate the tiny CP violating parameter ϵ_K. In K^0 mixing, the $W_1 W_2$ box-diagrams are proportional to the following CKM factor combinations

$$\lambda^{LR}_q \lambda^{RL}_{q'} = V^L_{qs} V^{R*}_{qd} V^R_{q's} V^{L*}_{q'd} \qquad (q, q' = u, c, t), \tag{19}$$

The condition for (c, c) loop to be CP conserving leads to

$$\alpha_1 - \alpha_2 - \beta_1 \simeq 0 \quad \text{or} \quad \beta'_R \simeq -\beta'_L \simeq -\beta_L. \tag{20}$$

As all the CP phases α_i and β_i are expected to be small quantities, we neglect the possibility of $\alpha_1 - \alpha_2 - \beta_1 \simeq \pi$. For the charged Higgs contribution, the situation is similar to the general 2HDM: although the Yukawa couplings are less constrained and in general complex, the dominant contribution is only proportional to the left-handed CKM matrix V^L in the same manner as in the SM. Thus the charged Higgs contribution to (c,c) loop remains real up to $\mathcal{O}(\lambda^5)$, with λ the Wolfenstein parameter.

Since both of the two contributions are nearly real, their interference is either constructive or destructive. It will be shown bellow that due to the different chiralities, the contribution from charged Higgs loop interferes always destructively with the W_1W_2 loop in the CP conserving case of Eq.(20). This provides a possibility of a nearly complete cancellation, which may greatly reduce the mass lower bounds for both W_2 and charged Higgs H^+.

The SM W_1W_1 loop diagram contribution to $K^0-\bar{K}^0$ mixing is described by the following effective Hamiltonian

$$\begin{aligned} H_{eff}^{W_1W_1} &= \frac{G_F^2 m_W^2}{16\pi^2}\left[(\lambda_c^{LL})^2\eta_{cc}S_0(x_c) + (\lambda_t^{LL})^2\eta_{tt}S_0(x_t) + 2\lambda_c^{LL}\lambda_t^{LL}\eta_{ct}S_0(x_c,x_t)\right] \\ &\quad \bar{d}\gamma^\mu(1-\gamma_5)s\otimes\bar{d}\gamma_\mu(1-\gamma_5)s + \text{H.c}, \end{aligned} \tag{21}$$

where m_W is the mass of the left-handed gauge boson and

$$S_0(x) = \frac{x}{(1-x)^2}\left[1 - \frac{11x}{4} + \frac{x^2}{4} - \frac{3x^2\ln x}{2(1-x)}\right], \tag{22}$$

$$S_0(x_c,x_t) = x_c x_t\left[-\frac{3}{4(1-x_c)(1-x_t)} + \frac{\ln x_c}{(x_c-x_t)(1-x_c)^2}\left(1 - 2x_c + \frac{x_c^2}{4}\right) + (x_c \leftrightarrow x_t)\right].$$

the Inami-Lim functions [48]. The CKM factors are defined as $\lambda_q^{LL} = V_{qs}^L V_{qd}^{L*}$. The matrix element is given by

$$\langle K^0|\bar{d}\gamma^\mu(1-\gamma_5)s\otimes\bar{d}\gamma_\mu(1-\gamma_5)s|\bar{K}^0\rangle = \frac{8}{3}f_K^2 m_K^2 B_K, \tag{23}$$

with the normalization $f_K = 159.8$ MeV . The values for QCD corrections are $\eta_{cc} \simeq 1.46$, $\eta_{ct} = 0.47 \pm 0.04$, and $\eta_{tt} = 0.5765 \pm 0.0065$ [49], and the bag parameter is $B_K = 0.86 \pm 0.06 \pm 0.14$[50].

The effective Hamiltonian for W_1W_2 loop and S(Goldstone)W_2 has been extensively investigated [51, 52, 53, 54, 55, 56, 57, 58] and reads

$$\begin{aligned} H_{eff}^{W_1W_2+SW_2} &= \frac{G_F^2 m_W^2}{8\pi^2}\beta\sum_{i,j}\lambda_i^{LR}\lambda_j^{RL}\sqrt{x_ix_j}\left[(4+x_ix_j\beta)\eta_1^{LR}I_1(x_i,x_j,\beta)\right. \\ &\quad \left.-(1+\beta)\eta_2^{LR}I_2(x_i,x_j,\beta)\right]\bar{d}(1-\gamma_5)s\otimes\bar{d}(1+\gamma_5)s + \text{H.c}, \end{aligned} \tag{24}$$

where $x_i = m_i^2/m_W^2$ and $\beta = m_W^2/M_2^2$. The two QCD correction coefficients are $\eta_1^{LR} = 1.4$ and $\eta_2^{LR} = 1.17$ for $\Lambda_{QCD} = 0.2$GeV[32]. The phases in CKM matrix elements are defined as $\lambda_q^{LR} = V_{qs}^L \left(V_{qd}^R\right)^*$ and $\lambda_q^{RL} = V_{qs}^R \left(V_{qd}^L\right)^*$ with the loop functions[57]

$$
\begin{aligned}
I_1(x_i, x_j, \beta) &= \frac{x_i \ln x_i}{(1-x_i)(1-x_i\beta)(x_i-x_j)} + (i \to j) - \frac{\beta \ln \beta}{(1-\beta)(1-x_i\beta)(1-x_j\beta)} \\
I_2(x_i, x_j, \beta) &= \frac{x_i^2 \ln x_i}{(1-x_i)(1-x_i\beta)(x_i-x_j)} + (i \to j) - \frac{\ln \beta}{(1-\beta)(1-x_i\beta)(1-x_j\beta)}. \qquad (25)
\end{aligned}
$$

In the limit of $x_i = x_j$ and $\beta \ll 1$, they reduce to

$$
I_1(x, \beta) \simeq \frac{1}{(1-x)} + \frac{\ln x}{(1-x)^2} - \beta \ln \beta, \qquad (26)
$$

$$
I_2(x, \beta) \simeq \frac{x}{1-x} + \frac{(2-x)x \ln x}{(1-x)^2} - \ln \beta. \qquad (27)
$$

Compared with the loop functions S_0 for W_1W_1 loop, the functions $I_1(x,\beta)$ and $I_2(x,\beta)$ have different mass dependencies: (i) For very small $x_c \ll 1$, the loop functions can be further simplified as $I_1(x_c, \beta) \simeq \ln x_c + 1$, and $I_2(x_c, \beta) \simeq -\ln \beta$. It is easy to see that the combination $4\eta_1^{LR} I_1(x_c, \beta) - \eta_2^{LR} I_2(x_c, \beta)$ is always *negative*. The sign of the amplitude is of crucial importance when there are multiple sources of contributions. The negative contribution may lead to cancellations with other amplitudes such as Higgs loops. (ii) The functions grow slowly with the internal quark masses. The typical values (for $M_2 = 1$ TeV) are $\beta x_c I_{1(2)}(x_c, \beta) = -1.3 \times 10^{-5}(9.2 \times 10^{-6})$ and $\beta x_t I_{1(2)}(x_t, \beta) = -4.3 \times 10^{-3}(0.077)$. Comparing with $S_0(x_c(x_t)) = 2.8 \times 10^{-4}(2.6)$, one sees that for $t-$quark loop, the loop functions $I_1(x_t, \beta)$ and $I_2(x_t, \beta)$ are much smaller than $S_0(x_t)$, which significantly suppresses the W_1W_2 loop contribution to B^0 mixing. While for $c-$quark loop, the W_1W_2 loop correction can be significant. Thus the main constraint for this model comes from neutral K meson system. The matrix element for the scalar operator is given by

$$
\langle K^0 | \bar{d}(1-\gamma_5)s \otimes \bar{d}(1+\gamma_5)s | \bar{K}^0 \rangle = \left[\frac{1}{3} + \frac{2m_K^2}{(m_s+m_d)^2} \right] f_K^2 m_K^2 B_K^S \qquad (28)
$$

with B_K^S the bag factor for scalar operator.

Since the $c-$quark mass dominates over the s, d quarks, only the first term in eq.(18) is considered which involves left-handed CKM matrix V^L. The dominant contributions are

$$
\begin{aligned}
H_{eff}^{H^\pm W_1 + H^\pm H^\pm} &= \frac{G_F^2}{16\pi^2} m_W^2 x_c y_c (\lambda_c^{LL})^2 \left[2\eta_{cc}^{HW} |\xi_c|^2 B_V^{HW}(y_c, y_W) + \frac{1}{4}\eta_{cc}^{HH} |\xi_c|^4 B_V^{HH}(y_c, y_W) \right] \\
&\cdot \bar{d}\gamma^\mu(1-\gamma_5)s \otimes \bar{d}\gamma_\mu(1-\gamma_5)s + \text{H.c}, \qquad (29)
\end{aligned}
$$

with $y_c = m_c^2/m_{H^\pm}^2$ and $y_W = m_W^2/m_{H^\pm}^2$. The loop functions are given by [59]

$$\begin{aligned} B_V^{HW}(y, y_W) &= \frac{y_W - \frac{1}{4}}{(1-y)(y-y_W)} + \frac{y_W - \frac{1}{4}y}{(1-y)^2(1-y_W)}\ln y + \frac{3}{4}\frac{y_W^2 \ln(y_W/y)}{(y_W-y)^2(1-y_W)}, \\ B_V^{HH}(y, y_W) &= \frac{1+y}{(1-y)^2} + \frac{2y}{(1-y)^3}\ln y. \end{aligned} \tag{30}$$

Note that the $H^\pm H^\pm$ loop is proportional to $|\xi_c^4|$, which significantly enhances the charged Higgs contribution at large ξ_c.

The contributions to the effective Hamiltonian can also arise from the flavor changing neutral current interactions via neutral Higgs exchange at tree level, which is denoted as $H_{eff}^{h^0}$. Summing up all the individual contributions, the total effective Hamiltonian is

$$H_{eff} = H_{eff}^{W_1 W_1} + H_{eff}^{W_1 W_2 + S W_2} + H_{eff}^{H^\pm W_1 + H^\pm H^\pm} + H_{eff}^{h^0}. \tag{31}$$

A. $K^0 - \bar{K}^0$ mass difference

The mass difference for K meson is simply given by

$$\Delta m_K \simeq 2\mathrm{Re}(M_{12}), \tag{32}$$

which is dominated by internal $c-$ quark loop for all loop diagrams, and can be calculated numerically. In Fig.1 the individual contribution from $W_1 W_2$ diagram is shown, the figure indicates a large negative M_{12} relative to the observed K^0 mass difference Δm_K at low $M_2 < 1$TeV, which reflects the difficulties to have a light W_2 bellow TeV in the one Higgs bi-doublet LR model. The situation is different when the charged Higgs contribution is taken into account. For the same set of CKM matrix element, the charged Higgs gives a positive contribution to M_{12}, which is comparable to the W_2 term for light charged Higgs mass with sufficiently large couplings. Taking only the dominant $H^\pm H^\pm$ loop contribution, we find that a cancellation between $W_1 W_2$ and $H^\pm H^\pm$ loop requires

$$\eta_{cc}^H x_c |\xi_c|^4 \frac{M_2^2}{m_H^2} \simeq -24\left[\eta_1^{LR}(\ln x_c + 1) + \frac{\eta_2^{LR}}{4}\ln\frac{m_W^2}{M_2^2}\right]\frac{m_K^2}{(m_s+m_d)^2}\frac{B_K^S}{B_K}. \tag{33}$$

A numerical calculation including all the contributions is shown in Fig.2. Numerically, for $m_{H^+} \sim 150$GeV and Yukawa coupling $\xi \sim 25$, the charged Higgs can compensate a opposite contribution from a light W_R at $M_2 \sim 600$ GeV. The large Higgs contribution relies on the fact that the $H^\pm W_1$ loop is proportional to $|\xi_c|^2$ and $H^\pm H^\pm$ loop proportional to $|\xi_c|^4$, which grow rapidly with $|\xi^c|$ increasing.

In Fig.3 we give total loop contributions from W_1W_2 loop, W_1H loop, HH loop together with the W_1W_1 loop in the SM. The mass of W_2 is set to 600GeV. One sees that in the range of 150GeV$< m_H^+ <$ 250GeV and $25 < \xi_c < 30$, the whole contribution can coincide with the experimental data of $\Delta m_K = (3.483 \pm 0.006) \times 10^{-15}$GeV. Generically, the needed charged Higgs mass m_H^+ grows with the mass of W_2. For a heavier W_2 at 1TeV and the same Yukawa coupling ξ_c the allowed value of m_H^+ is around $250 \sim 350$ GeV. The numerical result for this case is shown in Fig.4.

We now check the contributions from the flavor changing neutral current interactions via neutral Higgs exchanges at tree level. From eq.(18), we have

$$M_{12}^{h^0} \simeq 6.0 \times 10^{-16} \left(\frac{200\text{GeV}}{m_{h^0}}\right)^2 \frac{(\eta_{12}^d - \eta_{21}^{d*})^2}{(0.1)^2}, \tag{34}$$

where we have taken $m_d = 9$MeV and $m_s = 180$MeV. By requiring that the h^0 contribution can not excess the experimental data of Δm_K, one arrives at a upper bound of

$$\frac{\sqrt{\text{Re}[(\eta_{12}^d - \eta_{21}^{d*})^2]}}{m_{h^0}} \leq 8.3 \times 10^{-4}\text{GeV}^{-1}. \tag{35}$$

For m_{h^0} around 200 GeV, $\sqrt{\text{Re}[(\eta_{12}^d - \eta_{21}^{d*})^2]} \leq 0.16 \approx \mathcal{O}(0.1)$, which is in agreement with the approximate global U(1) family symmetry of $|\eta_{12}^d| \ll 1$. It will be shown below that a more stringent constraint can arise from the indirect CP violation ϵ_K. From that constraint, the contributions to the mass different from the flavor changing neutral Higgs interactions can be neglected.

B. Indirect CP violation ε_K

The indirect CP violation parameter ϵ_K arises from imaginary part of the mixing amplitude, dominated by internal (c, t) quarks and (t, t) quarks. With different CP phases, the interference among W_1W_1, W_1W_2 and $H^\pm H^\pm$ are rather complicated and the allowed parameter space is large.

The expression for indirect CP violation is given by

$$|\epsilon_K| \simeq \frac{1}{2\sqrt{2}} \left(\frac{\text{Im}M_{12}}{\text{Re}M_{12}} + 2\xi_0\right) \simeq \frac{\text{Im}M_{12}}{\sqrt{2}\Delta m_K}, \tag{36}$$

where ξ_0 is the weak phase of $K \to \pi\pi$ decay amplitude with isospin zero.

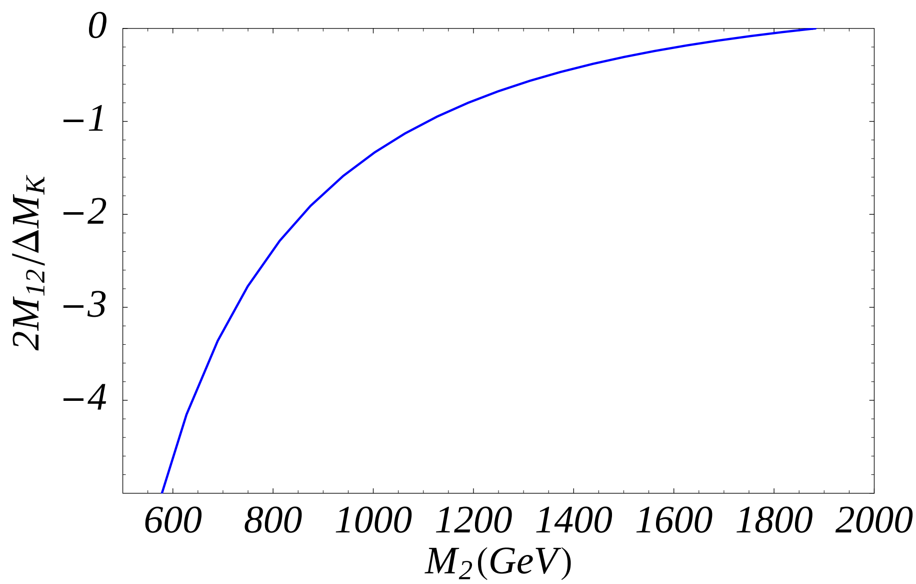

FIG. 1: Contribution to $2Re(M_{12})$ from right-handed W_R, normalized to the experimental data of Δm_K

Let us first examine the simplest case in which $|\xi_t|$ is tiny so that $H^\pm H^\pm$ loop is negligible. In this case the only extra contribution is from the $W_1 W_2$ loop. It is straight forward to see that once the (c,c) loop is set real as in Eq.(20), the (c,t) and (t,t) loop contributions become real as well because

$$\begin{aligned}\mathrm{Im}[\lambda_c^{LR}\lambda_t^{RL} + (L \to R)] &= -2|V_{cs}^L V_{cd}^L V_{ts}^{L*} V_{td}^{L*}| \cos(-\alpha_2 + \alpha_3 - \beta_2 + \phi) \\ &\quad \times \sin(\alpha_1 - \alpha_2 - \beta_1) \ , \end{aligned} \tag{37}$$

and

$$\mathrm{Im}(\lambda_t^{LR}\lambda_t^{RL}) = -|V_{ts}^L V_{td}^L|^2 \sin(\alpha_1 - \alpha_2 - \beta_1) \ . \tag{38}$$

where ϕ stands for $\arg(V_{cs}^L V_{cd}^L V_{ts}^{L*} V_{td}^{L*})$ and $\phi \simeq \beta_L' \simeq \beta_L$ in the Wolfenstein parametrization. Thus, there is no contribution to ϵ_K from $W_1 W_2$ loop.

In the case of non-negligible ξ_t, the charged Higgs contributes which involve the term proportional to $|\xi_t|^4$ and $|\xi_t \xi_c|^2$. The constraints are very similar to the general 2HDM case which has been analyzed in details in Ref.[25]. For a light charged Higgs mass $m_{H^+} = 150 \sim 300$ GeV, the typically allowed values are

$$|\xi_c| \sim \mathcal{O}(10) \quad \text{and} \quad |\xi_t| \sim \mathcal{O}(10^{-1}), \tag{39}$$

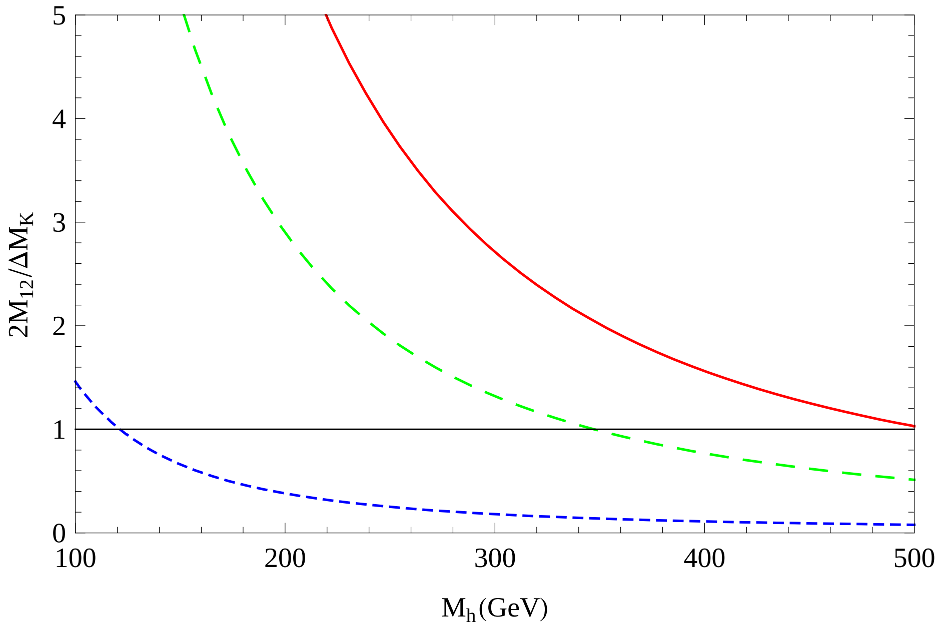

FIG. 2: Contribution to 2Re(M_{12}) from the lightest charge Higgs H^+. Three curves corresponds to ξ_c =30(solid), 25(dashed) and 15 (dotted)

.

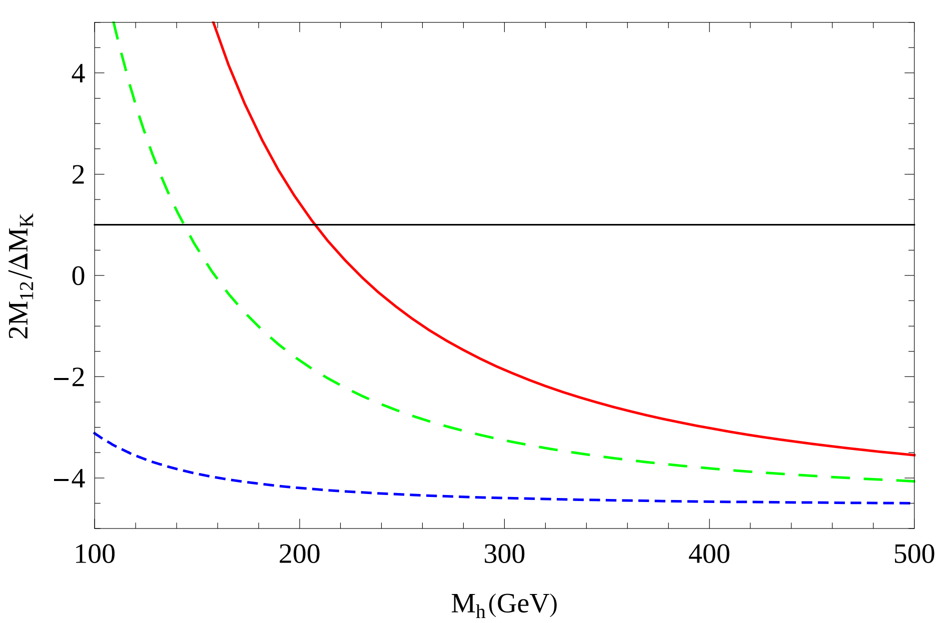

FIG. 3: Sum of all loop contributions, including the SM contribution to the $2M_{12}$ normalized to Δm_K with M_2 = 600 GeV. Three curves corresponds to $|\xi_c|$ =30(solid), 25(dashed) and 15 (dotted) respectively.

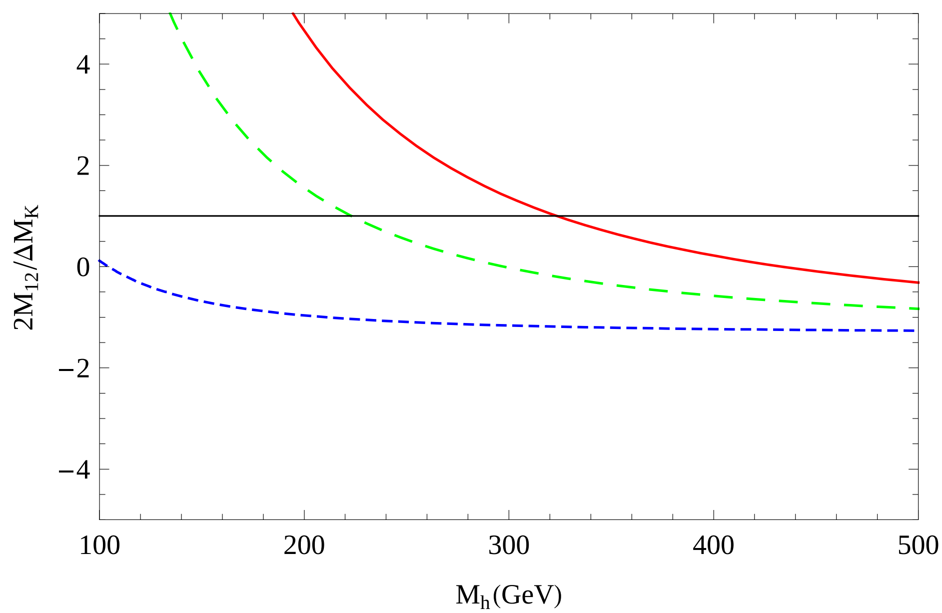

FIG. 4: The same as in Fig.3 with $M_2 = 1$TeV.

which is consistent with the previous discussion. As ξ_t is a free parameter, the constraints on ξ_c is not severe. Nevertheless, it indicates that a small $|\xi_t|$ is needed to meeting the experiments.

The Im$M_{12}^{h^0}$ contribution to the indirect CP violation ϵ_K from the flavor changing neutral current via neutral Higgs change at tree level can be significant

$$\epsilon_K^{h^0} \simeq 4.25 \times 10^{-4} \left(\frac{200\text{GeV}}{m_{h^0}}\right)^2 \frac{Im[(\eta_{12}^d - \eta_{21}^{d*})^2]}{(0.01)^2}. \tag{40}$$

From the requirement $|\epsilon_K^{h^0}| < \epsilon_K^{exp}$, we arrive at the constraint

$$\frac{|\text{Im}\left(\eta_{12}^d - \eta_{21}^{d*}\right)^2|^{1/2}}{m_{h^0}} < 6.9 \times 10^{-5}\text{GeV}^{-1}, \tag{41}$$

which is a more stringent constraint on the imaginary part of the Yukawa couplings. For $m_{h^0} = 200$GeV, one has $|\text{Im}\left(\eta_{12}^d - \eta_{21}^{d*}\right)^2|^{1/2} < 0.014$. It requires that either the off-diagonal coupling η_{12}^d should be very small or CP-violating phase must be tuned to be very small, or the neutral scalars must be very heavy, above 1 TeV. Here we shall consider the small off diagonal coupling via the mechanism of approximate U(1) family symmetry, which allows to have a light Higgs boson at the electroweak scale. In other words, the flavor changing neutral Higgs interactions in the two Higgs bi-doublet model can really be suppressed via such a mechanism, namely

$$|\eta_{12}^d| \leq 0.01, \quad \text{for} \quad m_{h^0} \sim 200 \text{ GeV}, \tag{42}$$

which is significantly different from the case in the one Higgs bi-doublet model[35, 36, 56, 57, 58] in which the off diagonal coupling is fixed to $(V^{L\dagger})_{2i} m_i^d (V^R)_{i1} v_1^* / (|v_1|^2 - |v_2|^2)$.

IV. NEUTRAL B MESON SYSTEM

In the previous section, we have illustrated that a light right-handed gauge boson can coincide with the K mixing data. In this section, we shall show that it is consistent with the B mixing measurements as well. Unlike the case in the Kaon system, the B meson mixing is dominated by internal $t-$quark loop. Due to the weaker dependence of loop functions $I_{1,2}(x,\beta)$ on quark masses. $I_1(x_t,\beta)$ and $I_2(x_t,\beta)$ from the $W_1 W_2$ loop are only a few percent of $S_0(x_t)$, which greatly suppresses it's phenomenological significance in B mixing and decays. The charged Higgs contribution is dominated by en extra parameter, the Yukawa coupling ξ_t, and is suppressed if $|\xi_t|$ is small.

In the first step let us take a close look at the $W_1 W_2$ contribution. The effective Hamiltonian for $\Delta B = 2$ process with $W_1 W_2$ loop is similar to the $\Delta S = 2$ case

$$
\begin{aligned}
H_{eff} \;\simeq\; & \frac{G_F^2 m_W^2}{8\pi^2} \lambda_t^{LR} \lambda_t^{RL} \beta x_t \left[(4 + x_t^2 \beta) I_1 \eta_1(x_t,\beta) - (1+\beta) I_2 \eta_2(x_t,\beta) \right] \\
& \bar{d}(1-\gamma_5) b \otimes \bar{d}(1+\gamma_5) b + \text{H.c.} \qquad (43)
\end{aligned}
$$

The QCD corrections at scale m_b are $\eta_1 \simeq 1.8$ and $\eta_2 \simeq 1.7$[57]. The matrix element is given by

$$
\begin{aligned}
M_{12}^{W_1 W_2} \;=\; & \frac{G_F^2 m_W^2}{8\pi^2} \lambda_t^{LR} \lambda_t^{RL} \beta x_t \left[4\eta_1 I_1(x_t,\beta) - \eta_2 I_2(x_t,\beta) \right] \left(\frac{m_B^2}{m_b^2} + \frac{1}{6} \right) f_B^2 m_B B_B^S \\
=\; & \frac{G_F^2 m_W^2}{8\pi^2} \left| V_{td}^L \right|^2 e^{-i(\beta_L + \beta_R)} \frac{m_t^2}{M_2^2} \left[4\eta_1 I_1(x_t, \frac{m_W^2}{M_2^2}) - \eta_2 I_2(x_t, \frac{m_W^2}{M_2^2}) \right] \qquad (44) \\
& \times \left(\frac{m_B^2}{m_b^2} + \frac{1}{6} \right) f_B^2 m_B B_B^S. \qquad (45)
\end{aligned}
$$

In the limit that all α s are vanishing, one has $\beta_R = -\beta_L$. The bag parameters from QCD sum rule gives $B_B^S(m_b)/B_B(m_b) = 1.2 \pm 0.2$[34]. and $f_B \sqrt{B_B} = 0.228 \pm 0.030 \pm 0.010$ GeV. The total contribution is given by

$$
\begin{aligned}
M_{12} \;=\; & M_{12}^{SM} + M_{12}^{W_1 W_2} = \frac{G_F^2 m_W^2}{6\pi^2} \left| V_{td}^L \right|^2 e^{-2i\beta_L} f_B^2 m_B B_B \{ \eta_B S_0(x_t) \\
& + \frac{3}{4} e^{-i(\beta_L - \beta_R)} \frac{m_t^2}{M_2^2} \left[4\eta_1 I_1(x_t, \frac{m_W^2}{M_2^2}) - \eta_2 I_2(x_t, \frac{m_W^2}{M_2^2}) \right] \left(\frac{m_B^2}{m_b^2} + \frac{1}{6} \right) \frac{B_B^S}{B_B} \} \qquad (46)
\end{aligned}
$$

with QCD correction factor $\eta_B = 0.551 \pm 0.007$. The neutral B meson mass difference is given by

$$\Delta m_B \simeq 2\,|M_{12}|\,. \tag{47}$$

The latest data give $\Delta m_B = (3.337 \pm 0.003) \times 10^{-13}$ GeV. The B^0 and B^0_s mixing are the most important for determining the CKM matrix elements V_{td} and V_{ts} in the SM. The SM global fit gives $|V^L_{td}| = (7.4 \pm 0.8) \times 10^{-3}$[60]. In the presence of new physics. The connection between Δm_B and the CKM matrix elements is in general complicated (see. e.g,[60]). The right-handed gauge boson will contribute to both mixing amplitude and phases.

With the pollution from W_2, the time-dependent decay $B \to J/\psi K_S$ only measures an effective phase angle which may differ from β_L. The expression for β_{eff} is

$$2\beta_{eff} = \mathrm{Im}\left(\frac{q}{p}\frac{\bar{A}}{A}\right) = \mathrm{Im}\sqrt{\frac{M^*_{12}}{M_{12}}} = \arg\left(M^*_{12}\right) \tag{48}$$

Using the measured experimental value of Δm_B and β_{eff} one can obtain the value of β_L as a function of β_R only. In the limit $m^2_W \ll M_2$, β_L is close to β_{eff} we have in a good approximation

$$\tan 2\beta_L \simeq \tan 2\beta_{eff}\left[1 - r\frac{\sin(\beta_R - \beta_{eff})}{2\sin 4\beta_{eff}}\right], \tag{49}$$

where r is the ratio between W_1W_1 and W_1W_2 box diagrams

$$r = \frac{3m^2_t(4\eta_1 I_1(x_t,\beta) - \eta_2 I_2(x_t,\beta))}{4M^2_2 S(x_t)\eta_B}\left(\frac{m^2_B}{m^2_b} + \frac{1}{6}\right)\frac{B^S_B}{B_B}. \tag{50}$$

The above express also lead to a bound on β_L expressions

$$\tan 2\beta_{eff}\left[1 - \frac{r}{2\sin 4\beta_{eff}}\right] \leq \tan 2\beta_L \leq \tan 2\beta_{eff}\left[1 + \frac{r}{2\sin 4\beta_{eff}}\right]. \tag{51}$$

In Fig.(5), we plot the β_L as a function of β_R with different right-handed gauge boson mass M_2. As mentioned before, the W_2 contribution to B mixing is rather limited. One sees that for a light W_2 around 600 GeV, and β_R varying from $-180°$ to $180°$, the modification to β_L is less than $2°$.

Once the β_L is obtained, one can evaluate the matrix element $|V^L_{td}|$ which given in Fig.(6) as a function of β_R. One sees again that the changes in $|V^L_{td}|$ is small for the whole range of β_R. Comparing with the global SM fit value of $|V^L_{td}|$ the modifications is within the 1σ error range. Thus the model can easily accommodate both the data of Δm_B and $\sin 2\beta_{J/\psi}$. For the left-right model with only one Higgs bi-doublet, since both β_L and β_R are calculable

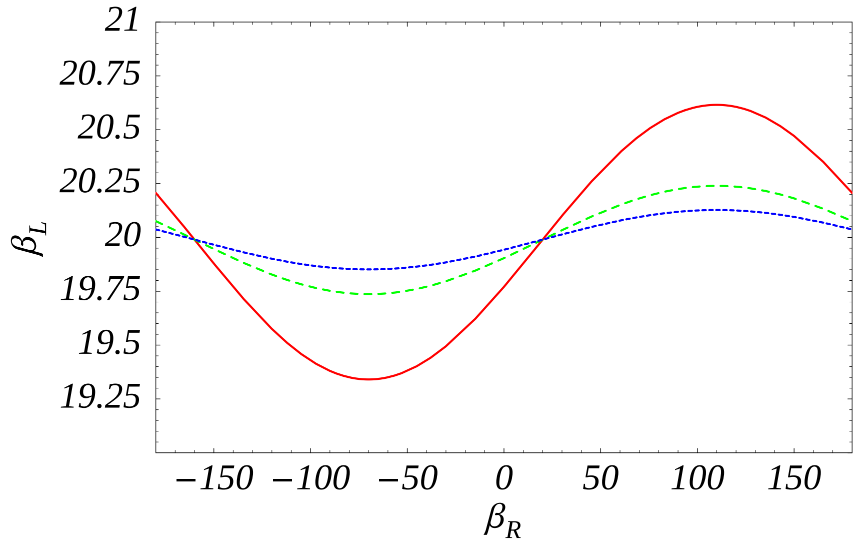

FIG. 5: Values of β_L as a function of β_R for different M_2. Three curves correspond to $M_2 = 500$ GeV (solid), 1000 GeV (dashed) and 1500 GeV(dotted) respectively.

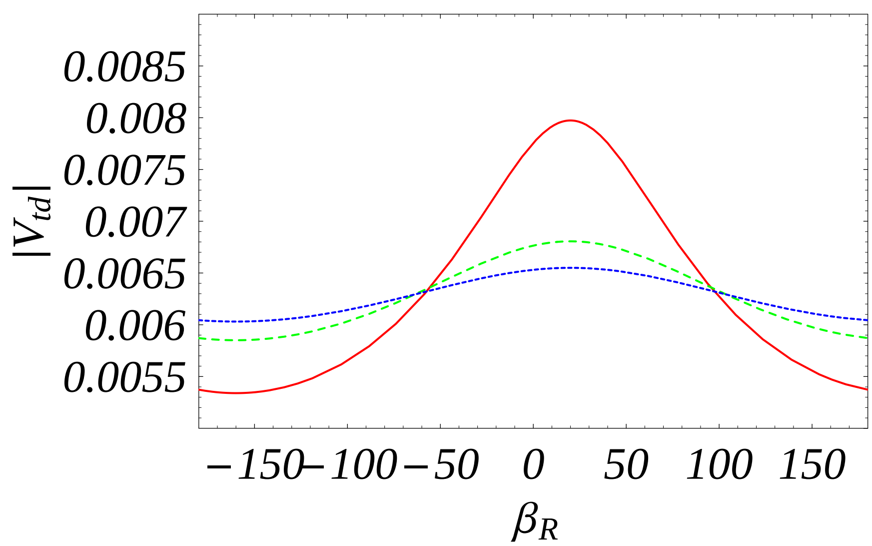

FIG. 6: Value of $|V_{td}|$as function of β_R. Three curves correspond to $M_2 = 500$ GeV (solid), 1000 GeV (dashed) and 1500 GeV(dotted) respectively.

quantities which depends only on the quark masses and ratios of VEVs, there is little room to meet the CP violation in both K and B system. Due to the suppression of CP phase form ϵ_K, the predicted $\sin 2\beta_{J/\psi}$ has to be small and can not excess 0.1.

The constraints from B system to charged Higgs couplings is quite similar to the general 2HDM. For $\xi \ll 1$, the box-diagram contribution can be safely neglected. The constraint

on the neutral Higgs FCNC couplings can be easily obtained from Eq.(18).

$$M_{12}^{h^0} \simeq 1.0 \times 10^{-12} \left(\frac{200\text{GeV}}{m_{h^0}}\right)^2 (\eta_{13}^d - \eta_{31}^{d*})^2. \tag{52}$$

Using the experimental data $\Delta m_B = 3.337 \times 10^{-3}$, one get a upper bound of

$$\frac{|\eta_{13}^d - \eta_{31}^{d*}|}{m_{h^0}} \leq 2.4 \times 10^{-3}\text{GeV}^{-1}. \tag{53}$$

For typical $m_{h^0} = 200$GeV, $|\eta_{13}^d - \eta_{31}^{d*}| < 0.41$.

Similarly, from the recently measured $\Delta m_{B_s} = 17.77 \pm 0.01 \pm 0.07\text{ps}^{-1}$[61], one can infer a bound for the couplings η_{23}^d and η_{32}^d

$$\frac{|\eta_{23}^d - \eta_{32}^{d*}|}{m_{h^0}} \leq 2.7 \times 10^{-3}\text{GeV}^{-1}. \tag{54}$$

For typical $m_{h^0} = 200$GeV, we get $|\eta_{13}^d - \eta_{31}^{d*}| < 0.54$. Both bounds satisfies the condition of $|\eta_{ij}^q| \ll 1$ from the approximate global $U(1)$ family symmetry.

V. CONCLUSIONS

In summary, motivated by natural spontaneous P and CP violation and the latest low energy experimental results, we have investigated a general left-right symmetric model with two Higgs bi-doublets. This simple extension evades the stringent constraints from K meson mixing, and lowers the allowed mass of right-handed gauge boson closing to the current direct experimental search bound ~ 600GeV. Through a negative interference with charged Higgs loop, which automatically occur when the charged Higgs is also light around electroweak scale with large Yukawa couplings. The FCNC can be suppressed by the mechanism of approximate global $U(1)$ family symmetry. We have illustrated that the off diagonal Yukawa couplings of $\mathcal{O}(10^{-2}) \sim \mathcal{O}(10^{-1})$ are consistent with all the constraints. This model has rich sources of CP violation which may show up in lower energy processes such as rare B decays and the new physics particles can be directly searched in upcoming LHC and future ILC experiments.

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